Models of inflation, supersymmetry breaking and observational constraints

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Abstract

We review the connection between inflationary models and observations and concentrate to describe models based on softly broken supersymmetry, in particular running mass models, and their predictions. We then present a fit of the spectral index of the curvature perturbation, assuming a flat ΛCDM cosmology.

1 Introduction

An epoch of inflationary expansion of the Universe is necessary to solve some of the problems of the standard Big Bang cosmology, like the large scale homogeneity and isotropy, flatness and unwanted relics problems. Moreover it has been also found that slow roll inflation can generate the small scale structure in the Universe through a primordial gaussian curvature perturbation, originated from the quantum fluctuations of the inflaton field.

In general a model of inflation consists in a scalar potential $V(\phi)$ for the inflaton field ϕ satisfying slow roll conditions[1] [1]:

$$\varepsilon = M_P^2 \left(\frac{V''}{V}\right)^2 \ll 1 \tag{1}$$

$$\eta = M_P^2 \left(\frac{V'}{V}\right)^2 \ll 1 \tag{2}$$

where the prime denotes derivative with respect to the field ϕ.

The point of contact between observation and models of inflation is the spectrum $\mathcal{P}_{\mathcal{R}}(k)$ of the curvature perturbation, which, in the slow roll approximation $3H\dot{\phi} \simeq -V'$, is given in terms of the inflaton potential $V(\phi)$

[1] As usual, in the following $M_P = 2.4 \times 10^{18}$GeV is the Planck mass, a is the scale factor and $H = \dot{a}/a$ is the Hubble parameter, and k/a is the wavenumber.

by

$$\frac{4}{25}\mathcal{P}_{\mathcal{R}}(k) = \frac{1}{75\pi^2 M_P^6}\frac{V^3}{V'^2}|_{k=aH}\,, \tag{3}$$

where the potential and its derivatives are evaluated at the epoch of horizon exit $k = aH$. To work out the value of ϕ at this epoch one uses the relation

$$\ln(k_{end}/k) \equiv N(k) = M_P^{-2}\int_{\phi_{end}}^{\phi}(V/V')d\phi\,, \tag{4}$$

where $N(k)$ is actually the number of e-folds from horizon exit of the scale k to the end of slow-roll inflation. At the scale explored by the COBE measurement of the cosmic microwave background (cmb) anisotropy, $N(k_{COBE})$ depends on the expansion of the Universe after inflation in the manner specified by:

$$N_{COBE} \simeq 60 - \ln(10^{16}\,\mathrm{GeV}/V^{1/4}) - \frac{1}{3}\ln(V^{1/4}/T_{\mathrm{reh}})\,. \tag{5}$$

In this expression, T_{reh} is the reheat temperature, and instant reheating is assumed.

Given the above relations, the observed large-scale normalization $\mathcal{P}_{\mathcal{R}}^{1/2} \simeq 10^{-5}$ provides a strong constraint on models of inflation. But here we are interested in the scale-dependence of the spectrum, defined by the, in general, scale-dependent spectral index n;

$$n(k) - 1 \equiv \frac{d\ln\mathcal{P}_{\mathcal{R}}(k)}{d\ln k}\,. \tag{6}$$

According to most inflationary models, n has negligible variation on cosmological scales so that $\mathcal{P}_{\mathcal{R}}(k) \propto k^{n-1}$, but we shall also discuss an interesting class of models giving a different scale-dependence.

From (3) and (4),

$$n - 1 \;=\; 2M_P^2(V''/V) - 3M_P^2(V'/V)^2\,, \tag{7}$$

and in almost all models of inflation, (7) is well approximated by

$$n - 1 = 2M_P^2(V''/V)\,. \tag{8}$$

We see that the spectral index measures the *shape* of the inflaton potential $V(\phi)$, being independent of its overall normalization. For this reason, it is a powerful discriminator between models of inflation.

2 Supersymmetric models of hybrid inflation

A very successful and interesting type of inflationary model is hybrid inflation, proposed by Linde [2]. In this case the potential depends on two fields, the inflaton and another field, which triggers a phase transition and the end of slow roll inflation and acquires a non vanishing v.e.v. at the end. The basic potential is of the form:

$$V = V_\phi(\phi) + \frac{1}{2}\left(\phi^2 - m_\psi^2\right)\psi^2 + ... \tag{9}$$

so, for $\phi \gg m_\psi$, the hybrid field ψ has vanishing v.e.v. and inflation can take place driven by V_ϕ, but when the inflaton becomes smaller, ψ is displaced from the origin and its v.e.v. stops slow roll inflation. Then a phase of oscillation of the two fields around the true minimum of the potential occurs and the preheating/reheating process.

We see then that the characteristic of hybrid inflation is that two different part of the potential are responsible for the inflationary phase and for the following stage of reheating. So the power spectrum of the curvature perturbations depends only on the inflationary potential V_ϕ and the critical value $\phi_c = m_\psi$, which corresponds to the end of slow roll inflation.

Supersymmetric versions of hybrid inflation have been studied by many authors (see [1] and references therein) and are based on the superpotential

$$W = \lambda\Phi(M^2 - \Psi\bar{\Psi}), \tag{10}$$

where capital Greek letter denote the superfields corresponding to the lower case Greek letter scalar field.

In the limit of global supersymmetry, such potential is perfectly flat with respect to the inflaton field and the small slope needed for slow roll is generated by supersymmetry breaking effects. Depending on the different mechanism considered, different scenarios are possible, as described in Table 1, together with the prediction for the spectral index and its derivative.

Note that the COBE normalization constraints the overall scale V_0, while the spectral index gives direct informations on the supersymmetry breaking effect which dominates the inflaton potential.

Table 1: Various models of hybrid inflation and their prediction for the spectral index as a function of $N = -\log(k/k_{end})$ and its derivative, assuming small field values. In the case of the simple linear term, $n-1$ is given by the full expression (7) since η vanishes, but note that in supergravity generally, other contribution to the spectral index coming from higher order terms are usually present and can be larger than the one listed here [3].

$V_\phi(\phi)/V_0$	$n-1$	$dn/d\log(k)$	Origin of the slope
$1-\frac{\lambda^2}{4\pi^2}\log\left(\frac{\sqrt{2}\lambda\phi}{Q}\right)$	$\frac{-1}{N+\frac{2\pi^2\phi_c^2}{\lambda^2 M_P^2}}$	$\frac{-1}{\left(N+\frac{2\pi^2\phi_c^2}{\lambda^2 M_P^2}\right)^2}$	1 loop for spont. broken susy
$1\pm\frac{1}{2}\frac{m_\phi^2}{V_0}\phi^2$	$\pm\frac{m_\phi^2 M_P^2}{V_0}$	0	Susy breaking mass
$1\pm\beta\frac{\phi}{M_P}$	$-3\beta^2$	0	Susy breaking linear term
$1\pm\frac{\phi^4}{M_P^4}$	$\frac{12}{\frac{M_P^2}{2\phi_c^2}\mp N}$	$\frac{\mp 12}{\left(\frac{M_P^2}{2\phi_c^2}\mp N\right)^2}$	Sugra quartic term

3 Running mass models

3.1 The potential

Let us consider now a specific type of models of hybrid inflation, the models with a running inflaton mass[4, 5, 6, 7, 8]. In these models, based on softly broken supersymmetry[2], one-loop corrections to the tree-level potential are taken into account by evaluating the inflaton mass-squared $m^2(\ln(Q))$ at the renormalization scale $Q \simeq \phi$, as long as ϕ greater than all other relevant scales.

$$V = V_0 + \frac{1}{2}m^2(\ln(Q))\phi^2 + \cdots . \tag{11}$$

The dependence on Q is only logarithmic due to supersymmetry, which ensures the cancelation of the quadratic divergences; anyway since supersymmetry is broken by the soft terms, a logarithmic divergence survives.

Over any small range of ϕ, it is a good approximation to take the running mass to be a linear function of $\ln\phi$ and then one reproduces explicitly the expression of the tree potential plus loop correction:

$$V = V_0 + \frac{1}{2}m^2(\ln(Q))\phi^2 - \frac{1}{2}c(\ln(Q))\frac{V_0}{M_P^2}\phi^2\ln(\phi/Q) + \cdots . \tag{12}$$

It has been shown [5] that the linear approximation is very good over the range of ϕ corresponding to horizon exit for scales between k_{COBE} and $8h^{-1}\mathrm{M}pc$. We shall want to estimate the reionization epoch, which involves a scale of order $k_{reion}^{-1} \sim 10^{-2}\mathrm{M}pc$ (enclosing the relevant mass of order $10^6\odot$). Since only a crude estimate of the reionization epoch is needed, we shall assume that the linear approximation is adequate down to this 'reionization scale'.

A potential of the type (12) gives rise to four different models of inflation, depending on the sign of c and the direction of rolling of the inflaton field, towards or away from the origin[3].

[2] We denote as spontaneously broken supersymmetry any theory where supersymmetry breaking preserves the super-trace relation $\mathcal{S}\mathrm{tr}\mathcal{M}^2 = 0$, while in softly broken supersymmetry one has $\mathcal{S}\mathrm{tr}\mathcal{M}^2 \neq 0$; e.g. the displacement of the inflaton field from the minimum of the potential causes a spontaneous breaking in the inflaton sector (which can give a soft breaking in the visible sector), while spontaneous or explicit supersymmetry breaking in another (hidden) sector is needed to have soft breaking in the inflaton sector. Note that in the first case, supersymmetry is restored once the inflaton field settles in the true minimum, while in the second case the breaking persists all the time, even if it can be affected by the dynamics during inflation.

[3] We follow the labeling introduced in [5].

The value of c is given by the well-known Renormalization Group Equation for the inflaton mass, and it depends on the gauge and Yukawa couplings, α and λ respectively, of the inflaton field:

$$\frac{V_0 c}{M_P^2} = -\frac{dm^2}{dt} = \frac{2C}{\pi}\alpha\widetilde{m}^2 - \frac{D}{16\pi^2}|\lambda|^2 m_{loop}^2 \,. \tag{13}$$

Here, C is a positive group-theoretic number of order 1, $\widetilde{m}$ is the supersymmetry breaking gaugino mass, D is a positive constant counting the number of scalar particles interacting with the inflaton through λ and m_{loop}^2 is their common susy breaking mass-squared. Negligible supersymmetry breaking trilinear coupling is assumed in the above expression.

We see that the running is then directly connected to the couplings of the inflaton field and the supersymmetry breaking mass spectrum of our model and so to complete our estimate of c, we need the gaugino or scalar mass. A very minimal and traditional hypothesis is that soft supersymmetry breaking is gravity-mediated and that the scale of susy breaking during inflation $V_0^{1/4}$ coincides with the scale of supersymmetry breaking in the present vacuum, $M_S \equiv \sqrt{F}$, where F is the auxiliary field responsible for spontaneous supersymmetry breaking in the hidden sector.

With gravity-mediated susy breaking, typical values of the masses are $\widetilde{m}^2 \sim |m_{loop}^2| \sim V_0/M_P^2$, which makes $|c|$ of order of the coupling strength α or $|\lambda|^2$.

At least in the case of dominance of the gauge coupling, one then expects a small, but non–negligible value for c

$$|c| \sim 10^{-1} \text{ to } 10^{-2} \,. \tag{14}$$

In special versions of gravity-mediated susy breaking, the masses could be much smaller, leading to $|c| \ll 1$. In that case, the mass would hardly run, and the spectral index would be practically scale-independent. With gauge-mediated susy breaking, the masses could be much bigger; this would not lead to a model of inflation (unless the coupling is very suppressed) because it would not satisfy the slow-roll requirement $|c| \lesssim 1$.

So we see that both from the theoretical and from the observational perspective, the interesting parameter region for the running mass models is given by (14).

3.2 The spectrum and the spectral index

Using (4) we find

$$s e^{c\Delta N(k)} = c\ln(\phi_*/\phi) \tag{15}$$

$$\Delta N(k) \equiv N_{COBE} - N(k) \equiv \ln(k/k_{COBE}), \tag{16}$$

where s is an integration constant, which absorbs the dependence on the end of inflation. The spectral index is then given, using (8), by

$$\frac{n(k)-1}{2} = se^{c\Delta N(k)} - c\,. \tag{17}$$

We see then that a relatively strong scale dependence arises for $n(k)$, depending on the magnitude of the constants s and c. To satisfy the slow-roll conditions (1), both c and s must be smaller than 1 in magnitude.

In order to evaluate the power spectrum later on (see (27) in the following section), we also need the variation of δ_H which comes from integrating this expression,

$$\frac{\delta_H(k)}{\delta_H(k_{COBE})} = \exp\left[\frac{s}{c}\left(e^{c\Delta N} - 1\right) - c\Delta N\right]. \tag{18}$$

We are mostly interested in cosmological scales between k_{COBE} and k_8, corresponding to $0 \lesssim \Delta N \lesssim 4$. In this range the scale-dependence of n is approximately linear (taking $|c| \lesssim 1$) and the variation $\Delta n \equiv n_8 - n_{COBE}$ is given approximately by

$$\Delta n \simeq 4\frac{dn}{d\ln k} \simeq 8sc\,. \tag{19}$$

4 Fit to a ΛCDM model

We will present here a global fit to the cosmological parameters according to the procedure described in full detail in [9]. Following [10] we will take into account also the recent measurements of the cmb anisotropy performed by the Boomerang [11] and Maxima-1 [12] balloon experiments.

The observational constraints on the cosmological parameters and on the spectral index have been studied by many authors, but in our analysis we will use a different treatment of reionization and we will compare the present data with the two parameters prediction for the scale dependence of the spectral index in running mass models.

Observations of various types indicate that we live in a low density Universe, which is at least approximately flat (see e.g. [11, 12]). In the interest of simplicity we therefore adopt the ΛCDM model, defined by the requirements that the Universe is exactly flat, and that the non-baryonic dark matter is cold with negligible interaction. Essentially exact flatness is predicted by inflation, unless one invokes a special kind of model, or special initial conditions. We shall constrain the parameters of the ΛCDM model, including

the spectral index, by performing a least-squares fit to key observational quantities.

4.1 The parameters

The ΛCDM model is defined by the spectrum $\mathcal{P}_{\mathcal{R}}(k)$ of the primordial curvature perturbation, and the four parameters that are needed to translate this spectrum into spectra for the matter density perturbation and the cmb anisotropy. The four parameters are the Hubble constant h (in units of $100\mathrm{k}ms^{-1}\mathrm{M}pc^{-1}$), the total matter density parameter Ω_0, the baryon density parameter Ω_b, and the reionization redshift z_R. As we shall describe, z_R is estimated by assuming that reionization occurs when some fixed fraction f of the matter collapses. Within the reasonable range $f \sim 10^{-4}$ to 1, the main results are insensitive to the precise value of f.

The spectrum is conveniently specified by its value at a scale explored by COBE, and the spectral index $n(k)$. We shall consider the usual case of a constant spectral index, and the case of running mass models where $n(k)$ is given by the two-parameter expression in (17). Since $\mathcal{P}_{\mathcal{R}}(k_{COBE})$ is determined very accurately by the COBE data we fix its value. Excluding z_R and $\mathcal{P}_{\mathcal{R}}(k_{COBE})$, the ΛCDM model is specified by four parameters in the case of a constant spectral index, or by five parameters in the case of running mass inflation models.

4.2 The data

As described in detail in [9, 10], we consider a sample of data constituted by seven observational quantities. Of these quantities, three are the cosmological quantities h, Ω_0, Ω_B, which we are also taking as free parameters. We will assume that, at least at some crude level, we can pretend that the errors are all random and uncorrelated, and perform a least squares fit.

The other data are taken so to sample different observable scales: one one side we use large scale structure observation, on the other measurements of the cmb anisotropy.

So, first of all we consider the rms density perturbation at $8h^{-1}$Mpc, σ_8, measured through the abundance of rich galaxy clusters at redshift $z = 0$ to a few [13] and the shape parameter Γ that specifies the slope of the galaxy correlation function on scales of order $1h^{-1}$ to $100h^{-1}\mathrm{M}pc$ [14, 15]

Secondly we take also two data from the latest measurements of the cmb anisotropy, i.e. the height of the first peak at $\ell \simeq 210 - 30$ and the ratio

between the heights of the second and first peak. We consider the average of the two experiment [11, 12] and include the calibration error in quadrature.

We fix the value of the COBE normalization, as described in detail in [9].

The adopted values and errors are given in Table 2 and 3, with the results of the fit for a constant and scale–dependent n.

4.3 Reionization

The effect of reionization on the cmb anisotropy is determined by the optical depth τ. We assume sudden, complete reionization at redshift z_R, so that the optical depth τ is given by [16]

$$\tau = 0.035\frac{\Omega_b}{\Omega_0}h\left(\sqrt{\Omega_0(1+z_R)^3+1-\Omega_0}-1\right). \tag{20}$$

In previous investigations, z_R has been regarded as a free parameter, usually fixed at zero or some other value. In this investigation, we instead estimate z_R, in terms of the parameters that we are varying plus assumed astrophysics. Indeed, it is usually supposed that reionization occurs at an early epoch, when some fraction f of the matter has collapsed into objects with mass very roughly $M = 10^6_\odot$. Estimates of f are in the range [17]

$$10^{-4.4} \lesssim f \lesssim 1\,. \tag{21}$$

In the case $f \ll 1$, the Press-Schechter approximation gives the estimate

$$1+z_R \simeq \frac{\sqrt{2}\sigma(M)}{\delta_c g(\Omega_0)}\,\mathrm{erfc}^{-1}(f) \qquad\qquad (f \ll 1)\,. \tag{22}$$

Here $\sigma(M)$ is the present, linearly evolved, rms density contrast with top-hat smoothing, and $\delta_c = 1.7$ is the overdensity required for gravitational collapse (g is the suppression factor of the linearly evolved density contrast at the present epoch, which does not apply at the epoch of reionization. See [9] for details). In the case $f \sim 1$, one can justify only the rough estimate

$$1+z_R \sim \frac{\sigma(M)}{g(\Omega_0)} \qquad\qquad (f \sim 1)\,, \tag{23}$$

not very different from the one that would be obtained by using $f = 1$ in (22).

In our fits, we fix f at different values in the above range, and find that the most important results are not very sensitive to f even though the corresponding values of z_R can be quite different.

4.4 The predicted peak height

The CMBfast package [18] gives C_ℓ, for given values of the parameters with n taken to be scale-independent. Following [19], we parameterize the CMBfast output at the first peak in the form

$$\widetilde{C}_{peak} = \widetilde{C}^{(0)}_{peak} \left(\frac{220}{10} \right)^{\nu/2} , \tag{24}$$

where $\widetilde{C}^{(0)}_{peak}$ is the value of $\widetilde{C}_{peak}$ at the reference value of the parameters, where $\nu = 0$, and

$$\nu \equiv a_n(n-1) + a_h \ln(h/0.65) + a_0 \ln(\Omega_0/0.35) + a_b h^2 (\Omega_b - 0.019) - 0.65 f(\tau)\tau \, . \tag{25}$$

The coefficients are $a_n = 0.88$, $a_h = -0.37$, $a_0 = -0.16$, $a_b = 5.4$, and $\widetilde{C}^{(0)}_{peak} = 77.5\mu\mathrm{K}$. The formula reproduces the CMBfast results within 10% for a 1-σ variation of the cosmological parameters, h, Ω_0 and Ω_b, and $n_{COBE} = 1.0 \pm 0.05$.

With the function $f(\tau)$ set equal to 1, the term -0.65τ is equivalent to multiplying $\widetilde{C}_{peak}$ by the usual factor $\exp(-\tau)$. We use the following formula, which was obtained by fitting the output of CMBfast, and is accurate to a few percent over the interesting range of τ;

$$f = 1 - 0.165\tau/(0.4 + \tau) \, . \tag{26}$$

For the running-mass model, we start with the above estimate for $n = 1$, and rescale it according to the scale dependence of (18), i.e.

$$\frac{\widetilde{C}_{peak}}{\sqrt{\widetilde{C}^{(n=1)}_{peak}}} = \frac{\delta_H(k(\ell, \Omega_0))}{\delta_H(k_{COBE}(\Omega_0))} \, . \tag{27}$$

In the case of constant n, this prescription corresponds to the previous one with $a_n = 0.91$, in good agreement with the output of CMBfast.

5 Results

5.1 Constant n

For the case of a constant spectral index our results are summarized in Table (2), together with the data points used.

The relative high value of the χ^2 is mainly due to the second peak contribution, but note that the discrepancy between the measured and fitted

Table 2: Fit of the ΛCDM model to presently available data, assuming reionization when a fraction $f = 10^{-2.2}$ of matter has collapsed (corresponding redshift at best fit is $z_R = 18$). The first two rows show the data points, while the result of the least-squares fit for the parameters is given in the lines three to five. All uncertainties are at the nominal 1-σ level. The total χ^2 is 6.3 with three degrees of freedom

	n	$\Omega_b h^2$	Ω_0	h	$\widetilde{\Gamma}$	$\widetilde{\sigma}_8$	$\sqrt{\widetilde{C}_\ell^{1st}}$	$\frac{\widetilde{C}_\ell^{2nd}}{\widetilde{C}_\ell^{1st}}$
data	—	0.019	0.35	0.65	0.23	0.56	$74.0\,\mu K$	0.38
err.	—	0.002	0.075	0.075	0.035	0.059	$5.0\,\mu K$	0.06
fit	0.99	0.021	0.38	0.62	0.19	0.56	$70.8\,\mu K$	0.49
err.	0.05	0.002	0.06	0.05	—	—	—	—
χ^2	—	0.9	0.2	0.2	1.3	0.002	0.4	3.3

values of the peak ratios is within the 2σ limit and that the best fit value of $\Omega_b h^2$ is also within 1σ from the Big Bang Nucleosynthesis measurement.

We show also in Figure 1, the comparison of the dependence of the best fit value of the spectral index on the the reionization redshift with the dependence on the reionization fraction f. Note that a strong correlation between n and z_R is present, while the correlation with f is much weaker, and in particular the lower bound on n is practically independent of f. Low values of z_R and therefore high f have lower χ^2.

5.2 Running mass models

We performed the same fit using also the scale dependent spectral index given in (17). In this case we obtain an allowed region in the plane s vs c, shown in Figure 2. In the same figure it is also displayed the most reasonable region of values for the parameters, from theoretical arguments (see [9] for a full discussion).

In Table 3 are again summarized the results of the fit in this case: we have fixed the value of c, so that the number of degrees of freedom is the same as before. The value of n in the table, refers to a scale corresponding to the COBE measurement. Since c is positive, smaller scales present larger spectral index $n > n_{COBE}$, e.g. $n(k_8) = 1.009$. Note that the value of χ^2 is

Table 3: Fit of the ΛCDM model to presently available data for the running mass models. We have fixed $c = 0.1$ and so the free parameters are $n_{COBE} = 1 + 2(s - c)$, and the next three quantities in the Table. Reionization is taken to occur when a fraction $f = 10^{-2.2}$ of matter has collapsed, as done previously (the corresponding redshift at best fit is $z_R = 21$.) Every quantity except n_{COBE} is a data point, with the value and uncertainty listed in the first two rows. The result of the least-squares fit is given in the lines three to five. All uncertainties are at the nominal 1-σ level. The total χ^2 is 8.4 with three degrees of freedom

	$n_{\rm COBE}$	$\Omega_b h^2$	Ω_0	h	$\widetilde{\Gamma}$	$\widetilde{\sigma}_8$	$\widetilde{C}_\ell^{1st}$	$\frac{\widetilde{C}_\ell^{2nd}}{\widetilde{C}_\ell^{1st}}$
data	—	0.019	0.35	0.65	0.23	0.56	74.0 μK	0.38
err.	—	0.002	0.075	0.075	0.035	0.059	5.0 μK	.06
fit	0.94	0.021	0.40	0.59	0.19	0.53	67.6 μK	0.49
err.	0.04	0.002	0.05	0.05	—	—	—	—
χ^2	—	0.9	0.4	0.6	1.2	0.2	1.6	3.5

in this case larger: the minimum lies on the $c = 0$ line and favours no scale–dependence. Note anyway that along the quasi degenerate $s - c$ direction, corresponding to small n_{COBE}, values of $|c| \leq 0.1 - 0.2$ are within the 70% CL contour.

6 Conclusion

We have presented a fit of the ΛCDM model to a global data set, assuming that a gaussian primordial curvature perturbation is the only one. We focused on the spectral index n, specifying the shape of the curvature perturbation, considering separately the case of a practically scale-independent spectral index, and the scale-dependent spectral index predicted by running mass inflation models. In contrast with other groups, we calculate the reionization epoch within the model on the assumption that it corresponds to the epoch when some fraction f of the matter collapses, the results being only mildly dependent on f in the reasonable range $f \gtrsim 10^{-4}$.

For the scale-independent case, we obtain a spectral index $n = 0.99 \pm 0.05$ for an intermediate value of f. The result for all f is given in Figure 1. We stress that the best fit result strongly depends on the procedure used to treat the reionization redshift z_R.

In the case of running mass models, the scale-dependent spectral index depends on parameters s and c, the latter being related to the inflaton coupling which produces the running. We have delineated the allowed region in the s-c plane. An unsuppressed coupling $c \sim 0.1$ is allowed by the data, leading to a noticeable scale-dependence of the spectral index. The fit with $c = 0.1$ is less good than with a scale-independent spectral index, but still acceptable.

Acknowledgments

I am very grateful and indebted to David H. Lyth, with whom this work has been done. I would like to thank prof. H.V. Klapdor-Kleingrothaus and the Organizing Committee of DARK2000 for their kind invitation to participate to this very interesting workshop and for partial financial support.

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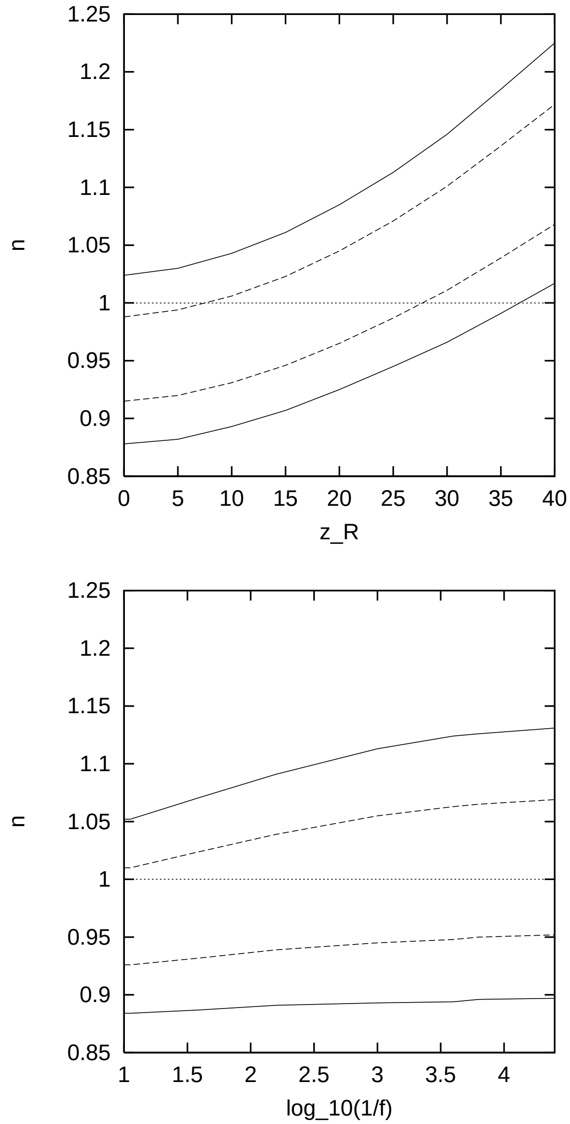

Figure 1: The plots show nominal 1- and 2-σ bounds on n. In (a) the fit is performed fixing the reionization epoch z_R, while in (b) is fixed instead the fraction f of matter which has collapsed at the epoch of reionization (the corresponding reionization redshifts, at best fit, are in the range 10 to 26)

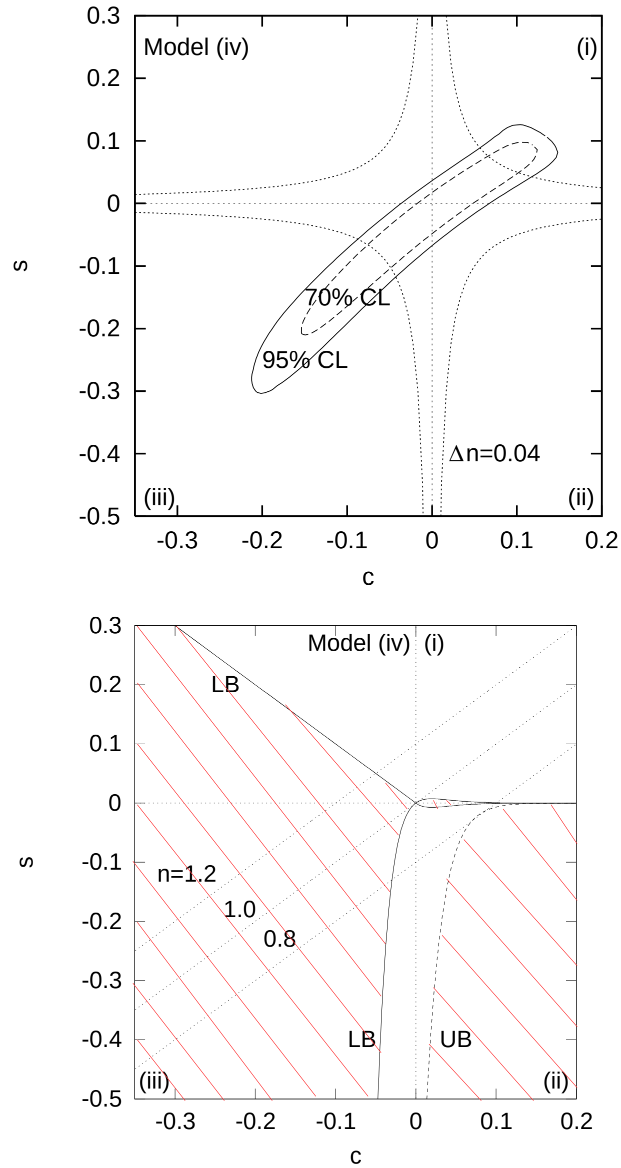

Figure 2: The parameter space for the running mass model. As discussed in [10] the model comes in four versions, corresponding to the four quadrants of the parameter space. In panel (a), we show the region allowed by observation, in the case that reionization occurs when $f \sim 1$; the scale-dependence of the prediction for n is also displayed in this panel by the branches of the hyperbola $8sc = \Delta n \equiv n_8 - n_{COBE}$, for the reference value $\Delta n = 0.04$. In panel (b), the straight lines correspond to $n_{COBE} = 1.2, 1.0$ and 0.8, and the shaded region is disfavoured on theoretical grounds